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## GRAVITY-ASSISTED EMERGENT HIGGS MECHANISM IN THE POST-INFLATIONARY EPOCH \*

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We consider a non-standard model of gravity coupled to a neutral scalar “inflaton” as well as to  $SU(2) \times U(1)$  iso-doublet scalar with positive mass squared and without self-interaction, and to  $SU(2) \times U(1)$  gauge fields. The principal new ingredient is employing two alternative non-Riemannian space-time volume-forms (covariant integration measure densities) independent of the metric. The latter have a remarkable impact – although not introducing any additional propagating degrees of freedom, their dynamics triggers a series of important features: appearance of infinitely large flat regions of the effective “inflaton” potential as well as dynamical generation of Higgs-like spontaneous symmetry breaking effective potential for the  $SU(2) \times U(1)$  iso-doublet scalar.

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### 1. Introduction

In a remarkable paper from 1986<sup>1</sup> J. Bekenstein proposed the intriguing idea about a gravity-assisted spontaneous symmetry breaking of electro-weak (Higgs) type without invoking unnatural (according to Bekenstein’s opinion) ingredients like negative mass squared and a quartic self-interaction for the Higgs field. By considering a model of gravity interacting with a standard Klein-Gordon field (with small positive mass squared and without self-interaction) coupled conformally to the scalar curvature he managed to obtain a prototype of dynamically induced Higgs-like spontaneous symmetry breaking scalar potential. A similar approach was further worked out in Ref. 2.

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Motivated by Bekenstein's idea, in the present essay we will consider a non-standard model of gravity coupled to a neutral scalar “inflaton”  $\varphi$  as well as to a  $U(1)$ -charged  $SU(2)$  iso-doublet scalar field  $\sigma$  with a standard positive mass squared and no self-interaction, as well as to  $SU(2) \times U(1)$  gauge fields. The essential non-standard feature of this model is employing non-Riemannian space-time volume forms – alternative generally covariant integration measure densities defined in terms of auxiliary antisymmetric tensor gauge fields independent of the pertinent Riemannian metric<sup>a</sup>. Although being almost pure-gauge degrees of freedom (see the remark in Section 2 below), the non-Riemannian space-time volume forms trigger a series of important features unavailable in ordinary gravity-matter models with the standard Riemannian volume-form (given by the square-root of the determinant of the Riemannian metric): (i) The “inflaton”  $\varphi$  develops a remarkable effective scalar potential in the Einstein frame possessing an infinitely large flat region for large negative  $\varphi$  describing the “early” universe evolution; (ii) In the absence of the  $SU(2) \times U(1)$  iso-doublet scalar field, the “inflaton” effective potential has another infinitely large flat region for large positive  $\varphi$  describing the “late” post-inflationary (dark energy dominated) universe; (iii) Inclusion of the  $SU(2) \times U(1)$  iso-doublet scalar field  $\sigma$  introduces a drastic change in the total effective scalar potential in the post-inflationary universe – the effective potential as a function of  $\sigma$  acquires exactly the electro-weak Higgs-type spontaneous symmetry breaking form.

## 2. Gravity-Gauge-Field-Matter Model With Two Independent Non-Riemannian Volume-Forms

Our starting point is the following non-standard gravity-gauge-field-matter system with an action of the general form involving two independent non-Riemannian volume-forms generalizing the model studied in Ref. 6 and whose gauge-field-matter part has an internal  $SU(2) \times U(1)$  gauge symmetry (for simplicity we will use units where the Newton constant is taken as  $G_{\text{Newton}} = 1/16\pi$ ):

$$S = \int d^4x \Phi_1(A) \left[ R - 2\Lambda_0 \frac{\Phi_1(A)}{\sqrt{-g}} + L^{(1)} \right] + \int d^4x \Phi_2(B) \left[ L^{(2)} + \frac{\Phi(H)}{\sqrt{-g}} \right]. \quad (1)$$

In (1) and in the sequel the following notations are used:

- $\Phi_1(A)$  and  $\Phi_2(B)$  are two independent non-Riemannian volume-forms, i.e., alternative generally covariant integration measure densities on the underlying space-time manifold:

$$\Phi_1(A) = \frac{1}{3!} \varepsilon^{\mu\nu\kappa\lambda} \partial_\mu A_{\nu\kappa\lambda} \quad , \quad \Phi_2(B) = \frac{1}{3!} \varepsilon^{\mu\nu\kappa\lambda} \partial_\mu B_{\nu\kappa\lambda} \quad , \quad (2)$$

<sup>a</sup>The method of non-Riemannian space-time volume forms, originally proposed in Refs. 3, 4 has been employed in a broad variety of gravity-matter, supergravity and strings/branes models to provide plausible solutions for dynamical generation of cosmological constant, the supersymmetric Higgs effect,<sup>5</sup> unified description of early universe inflation and present day dark energy,<sup>6</sup> unified description of dark energy and dark matter as different manifestations of a single entity<sup>7,8</sup>

defined in terms of field-strengths of two auxiliary 3-index antisymmetric tensor gauge fields.  $\Phi_{1,2}$  take over the role of the standard Riemannian integration measure density  $\sqrt{-g} \equiv \sqrt{-\det \|g_{\mu\nu}\|}$  in terms of the space-time metric  $g_{\mu\nu}$ .

- $R = g^{\mu\nu} R_{\mu\nu}(\Gamma)$  and  $R_{\mu\nu}(\Gamma)$  are the scalar curvature and the Ricci tensor in the first-order (Palatini) formalism, where the affine connection  $\Gamma_{\nu\lambda}^\mu$  is *a priori* independent of the metric  $g_{\mu\nu}$ .  $\Lambda_0$  is a small parameter later to be identified with the present epoch small observable cosmological constant.
- $L^{(1)}$  is the sum of two scalar field Lagrangians:

$$L^{(1)} = -\frac{1}{2}g^{\mu\nu}\partial_\mu\varphi\partial_\nu\varphi - V_1(\varphi) - g^{\mu\nu}(\nabla_\mu\sigma_a)^*\nabla_\nu\sigma_a - V_0(\sigma), \quad (3)$$

where  $\varphi$  denotes the neutral “inflaton” inert under  $SU(2) \times U(1)$  and  $\sigma \equiv (\sigma_a)$  is a complex  $SU(2) \times U(1)$  iso-doublet scalar field with the isospinor index  $a = +, 0$  indicating the corresponding  $U(1)$  charge. The gauge-covariant derivative in (3) acting on the iso-doublet scalar  $\sigma$  reads:

$$\nabla_\mu\sigma = \left(\partial_\mu - \frac{i}{2}\tau_A\mathcal{A}_\mu^A - \frac{i}{2}\mathcal{B}_\mu\right)\sigma, \quad (4)$$

with  $\frac{1}{2}\tau_A$  ( $\tau_A$  – Pauli matrices,  $A = 1, 2, 3$ ) indicating the  $SU(2)$  generators and  $\mathcal{A}_\mu^A$  ( $A = 1, 2, 3$ ) and  $\mathcal{B}_\mu$  denoting the corresponding  $SU(2)$  and  $U(1)$  gauge fields. The pertinent scalar field potentials are:

$$V_1(\varphi) = f_1 \exp\{-\alpha\varphi\} \quad , \quad V_0(\sigma) = m_0^2 \sigma_a^* \sigma_a, \quad (5)$$

where  $\alpha, f_1$  are dimensionful positive parameters, and  $V_0(\sigma)$  is just the standard mass term for the iso-doublet  $\sigma_a$  with positive mass squared.

- $L^{(2)}$  is the sum of a second “inflaton” Lagrangian plus the canonical Lagrangians for the  $SU(2)$  and  $U(1)$  gauge fields  $\mathcal{A}_\mu^A, \mathcal{B}_\mu$ :

$$L^{(2)} = -\frac{b}{2}e^{-\alpha\varphi}g^{\mu\nu}\partial_\mu\varphi\partial_\nu\varphi + U(\varphi) - \frac{1}{4g^2}F^2(\mathcal{A}) - \frac{1}{4g'^2}F^2(\mathcal{B}), \quad (6)$$

where:

$$U(\varphi) = f_2 \exp\{-2\alpha\varphi\}, \quad (7)$$

with  $f_2$  another dimensionful positive parameter, whereas  $b$  is a dimensionless one, and (all indices  $A, B, C = (1, 2, 3)$ ):

$$F^2(\mathcal{A}) \equiv F_{\mu\nu}^A(\mathcal{A})F_{\kappa\lambda}^A(\mathcal{A})g^{\mu\kappa}g^{\nu\lambda}, \quad F^2(\mathcal{B}) \equiv F_{\mu\nu}(\mathcal{B})F_{\kappa\lambda}(\mathcal{B})g^{\mu\kappa}g^{\nu\lambda}, \quad (8)$$

$$F_{\mu\nu}^A(\mathcal{A}) = \partial_\mu\mathcal{A}_\nu^A - \partial_\nu\mathcal{A}_\mu^A + \epsilon^{ABC}\mathcal{A}_\mu^B\mathcal{A}_\nu^C, \quad F_{\mu\nu}(\mathcal{B}) = \partial_\mu\mathcal{B}_\nu - \partial_\nu\mathcal{B}_\mu. \quad (9)$$

- $\Phi(H)$  indicates the dual field strength of a third auxiliary 3-index antisymmetric tensor gauge field:

$$\Phi(H) = \frac{1}{3!}\epsilon^{\mu\nu\kappa\lambda}\partial_\mu H_{\nu\kappa\lambda}, \quad (10)$$

whose presence is crucial for non-triviality of the model.

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**Remark.** The systematic canonical Hamiltonian treatment of (1) (see *Appendix*) and, more generally, of any gravity-matter models built with non-Riemannian volume-forms<sup>9</sup> shows that the auxiliary 3-index tensor gauge fields  $A_{\mu\nu\kappa}$ ,  $B_{\mu\nu\kappa}$ ,  $H_{\mu\nu\kappa}$  are almost pure-gauge, i.e., they do not introduce any new propagating field degrees of freedom except for a few free integration constants in their respective equations of motion (see (14)-(15) below).

Let us note that the requirement for invariance (with the exception of the mass term  $V_0(\sigma)$  of the iso-doublet  $\sigma_a$ ) under the following global Weyl-scale symmetry:

$$g_{\mu\nu} \rightarrow \lambda g_{\mu\nu} \ , \ \varphi \rightarrow \varphi + \frac{1}{\alpha} \ln \lambda \ , \ A_{\mu\nu\kappa} \rightarrow \lambda A_{\mu\nu\kappa} \ , \ B_{\mu\nu\kappa} \rightarrow \lambda^2 B_{\mu\nu\kappa} \ , \quad (11)$$

$$\Gamma_{\nu\lambda}^\mu \ , \ H_{\mu\nu\kappa} \ , \ \sigma_a \ , \ \mathcal{A}_\mu^A \ , \ \mathcal{B}_\mu - \text{inert} \ ,$$

uniquely fixes the structure of the non-Riemannian-measure gravity-gauge-field-matter action (1). In particular, for the same reason we have multiplied by an appropriate exponential factor the “inflaton” kinetic term in  $L^{(2)}$  (6).

Following Ref. 6 we will now derive the effective Einstein-frame form of the dynamics described by (1).

Variation of the action (1) w.r.t. affine connection  $\Gamma_{\nu\lambda}^\mu$ :

$$\int d^4x \sqrt{-g} g^{\mu\nu} \left( \frac{\Phi_1}{\sqrt{-g}} \right) (\nabla_\kappa \delta \Gamma_{\mu\nu}^\kappa - \nabla_\mu \delta \Gamma_{\kappa\nu}^\kappa) = 0 \quad (12)$$

shows that  $\Gamma_{\nu\lambda}^\mu$  becomes a Levi-Civita connection  $\Gamma_{\nu\lambda}^\mu = \Gamma_{\nu\lambda}^\mu(\bar{g}) = \frac{1}{2} \bar{g}^{\mu\kappa} (\partial_\nu \bar{g}_{\lambda\kappa} + \partial_\lambda \bar{g}_{\nu\kappa} - \partial_\kappa \bar{g}_{\nu\lambda})$  w.r.t. to the Weyl-rescaled metric  $\bar{g}_{\mu\nu}$ :

$$\bar{g}_{\mu\nu} = \chi_1 g_{\mu\nu} \ , \ \chi_1 \equiv \frac{\Phi_1(A)}{\sqrt{-g}} \ . \quad (13)$$

Variation of (1) w.r.t. auxiliary tensor gauge fields  $A_{\mu\nu\lambda}$ ,  $B_{\mu\nu\lambda}$  and  $H_{\mu\nu\lambda}$  yields the equations (using the short-hand notation  $\chi_1$  from (13)):

$$\partial_\mu \left[ R + L^{(1)} - 4\Lambda_0 \chi_1 \right] = 0 \ , \ \partial_\mu \left[ L^{(2)} + \frac{\Phi(H)}{\sqrt{-g}} \right] = 0 \ , \ \partial_\mu \left( \frac{\Phi_2(B)}{\sqrt{-g}} \right) = 0 \ , \quad (14)$$

whose solutions read:

$$\frac{\Phi_2(B)}{\sqrt{-g}} \equiv \chi_2 = \text{const} \ , \ R + L^{(1)} - 4\Lambda_0 \chi_1 = -M_1 = \text{const} \ ,$$

$$L^{(2)} + \frac{\Phi(H)}{\sqrt{-g}} = -M_2 = \text{const} \ . \quad (15)$$

Here  $M_1$  and  $M_2$  are arbitrary dimensionful and  $\chi_2$  arbitrary dimensionless integration constants.

Now, varying (1) w.r.t.  $g_{\mu\nu}$  and using relations (15) we obtain:

$$\chi_1 \left[ R_{\mu\nu} + \frac{1}{2} \left( g_{\mu\nu} L^{(1)} - T_{\mu\nu}^{(1)} \right) - \Lambda_0 \chi_1 g_{\mu\nu} \right] - \frac{1}{2} \chi_2 \left[ T_{\mu\nu}^{(2)} + g_{\mu\nu} M_2 \right] = 0 \ , \quad (16)$$

where  $\chi_1$  and  $\chi_2$  are the same as in (13) and (15), respectively, and  $T_{\mu\nu}^{(1,2)}$  are the energy-momentum tensors of the matter-gauge-field Lagrangians (3)-(6) with the standard definitions:  $T_{\mu\nu}^{(1,2)} = g_{\mu\nu} L^{(1,2)} - 2\partial L^{(1,2)}/\partial g^{\mu\nu}$ .

Taking the trace of Eqs.(16) and using again the second relation (15) we solve for the scale factor  $\chi_1$  (13):

$$\chi_1 = 2\chi_2 \frac{T^{(2)}/4 + M_2}{L^{(1)} - T^{(1)}/2 - M_1} \quad , \quad T^{(1,2)} = g^{\mu\nu} T_{\mu\nu}^{(1,2)} . \quad (17)$$

Using the second relation (15), the gravity equations (16) can be put in the Einstein-like form:

$$R_{\mu\nu} - \frac{1}{2}g_{\mu\nu}R = \frac{1}{2}g_{\mu\nu} \left( L^{(1)} + M_1 - 2\Lambda_0\chi_1 \right) + \frac{1}{2} \left( T_{\mu\nu}^{(1)} - g_{\mu\nu}L^{(1)} \right) + \frac{\chi_2}{2\chi_1} \left[ T_{\mu\nu}^{(2)} + g_{\mu\nu}M_2 \right] , \quad (18)$$

Now, using expression (17) for  $\chi_1$  and following the same steps as in Ref. 6 we can bring Eqs.(18) into the standard form of Einstein equations for the rescaled metric  $\bar{g}_{\mu\nu}$  (13), i.e., the Einstein-frame gravity equations:

$$R_{\mu\nu}(\bar{g}) - \frac{1}{2}\bar{g}_{\mu\nu}R(\bar{g}) = \frac{1}{2}T_{\mu\nu}^{\text{eff}} \quad , \quad T_{\mu\nu}^{\text{eff}} = g_{\mu\nu}L_{\text{eff}} - 2\frac{\partial}{\partial g^{\mu\nu}}L_{\text{eff}} , \quad (19)$$

with effective energy-momentum tensor corresponding to the following effective Einstein-frame matter Lagrangian:

$$L_{\text{eff}} = A(\varphi, \sigma)X + B(\varphi)X^2 - U_{\text{eff}}(\varphi, \sigma) - \bar{g}^{\mu\nu}(\nabla_\mu\sigma_a)^*\nabla_\nu\sigma_a - \frac{\chi_2}{4g^2}\bar{F}^2(\mathcal{A}) - \frac{\chi_2}{4g'^2}\bar{F}^2(\mathcal{B}) . \quad (20)$$

In (20) the following notations are used:

- $X$  is the standard short-hand notations for the “inflaton” kinetic term:

$$X \equiv -\frac{1}{2}\bar{g}^{\mu\nu}\partial_\mu\varphi\partial_\nu\varphi . \quad (21)$$

- The coefficient functions in the “k-essence”-type<sup>10–13</sup> expression (first two terms on the r.h.s. of (20) with *non-linear* functional dependence on  $X$  (21)) are given by:

$$\begin{aligned} A(\varphi, \sigma) &\equiv 1 + \frac{1}{2}be^{-\alpha\varphi} \frac{V_1(\varphi) + V_0(\sigma) - M_1}{U(\varphi) + M_2} \\ &= 1 + \frac{1}{2}be^{-\alpha\varphi} \frac{f_1e^{-\alpha\varphi} + m_0^2\sigma_a^*\sigma_a - M_1}{f_2e^{-2\alpha\varphi} + M_2} ; \end{aligned} \quad (22)$$

and

$$B(\varphi) \equiv -\frac{1}{4} \frac{\chi_2 b^2 e^{-2\alpha\varphi}}{U(\varphi) + M_2} = -\frac{1}{4} \frac{\chi_2 b^2 e^{-2\alpha\varphi}}{f_2 e^{-2\alpha\varphi} + M_2} . \quad (23)$$

- The full effective scalar field potential reads:

$$\begin{aligned} U_{\text{eff}}(\varphi, \sigma) &\equiv \frac{\left[ V_1(\varphi) + V_0(\sigma) - M_1 \right]^2}{4\chi_2 \left[ U(\varphi) + M_2 \right]} + 2\Lambda_0 \\ &= \frac{\left( f_1 e^{-\alpha\varphi} + m_0^2 \sigma_a^* \sigma_a - M_1 \right)^2}{4\chi_2 \left[ f_2 e^{-2\alpha\varphi} + M_2 \right]} + 2\Lambda_0 , \end{aligned} \quad (24)$$

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where in(22)-(24) the explicit forms of  $V_1(\varphi)$ ,  $V_0(\sigma)$  and  $U(\varphi)$  (3)-(6) are inserted.

- In the last line of (20) the Einstein-frame versions of the gauge-field Lagrangians are used, where:

$$\bar{F}^2(\mathcal{A}) \equiv F_{\mu\nu}^A(\mathcal{A})F_{\kappa\lambda}^A(\mathcal{A})\bar{g}^{\mu\kappa}\bar{g}^{\nu\lambda} \quad , \quad \bar{F}^2(\mathcal{B}) \equiv F_{\mu\nu}(\mathcal{B})F_{\kappa\lambda}(\mathcal{B})\bar{g}^{\mu\kappa}\bar{g}^{\nu\lambda} . \quad (25)$$

### 3. Effective Scalar Potential – Infinitely Large Flat Regions and Emergent Higgs-like Behaviour

A remarkable feature of the effective Einstein-frame scalar potential (24) is that it possesses an infinitely large flat region for large negative  $\varphi$  and is independent of  $\sigma$  there:

$$U_{\text{eff}}(\varphi, \sigma) \simeq U_{(-)} \equiv \frac{f_1^2}{4\chi_2 f_2} + 2\Lambda_0 , \quad (26)$$

$$A(\varphi, \sigma) \simeq A_{(-)} \equiv 1 + \frac{1}{2}b\frac{f_1}{f_2} \quad , \quad B(\varphi) \simeq B_{(-)} \equiv -\chi_2\frac{b^2}{4f_2} .$$

Thus, for large negative values of the “inflaton” the effective matter Lagrangian (20) reads:

$$L_{\text{eff}}^{(-)} = A_{(-)}X + B_{(-)}X^2 - U_{(-)} + L[\sigma, \mathcal{A}, \mathcal{B}] , \quad (27)$$

$$L[\sigma, \mathcal{A}, \mathcal{B}] \equiv -\bar{g}^{\mu\nu}(\nabla_\mu\sigma_a)^*\nabla_\nu\sigma_a - \frac{\chi_2}{4g^2}\bar{F}^2(\mathcal{A}) - \frac{\chi_2}{4g'^2}\bar{F}^2(\mathcal{B}) , \quad (28)$$

i.e., the  $SU(2) \times U(1)$  iso-doublet scalar  $\sigma$  becomes massless.

For large positive values of  $\varphi$  we obtain  $A(\varphi, \sigma) \simeq A_{(+)} = 1$  ,  $B(\varphi) \simeq B_{(+)} = 0$ , and:

$$U_{\text{eff}}(\varphi, \sigma) \simeq U_{(+)}(\sigma) \equiv \frac{1}{4\chi_2 M_2} \left( m_0^2 \sigma_a^* \sigma_a - M_1 \right)^2 + 2\Lambda_0 , \quad (29)$$

so that the effective matter Lagrangian (20) becomes (using notation (28)):

$$L_{\text{eff}}^{(+)} = X - U_{(+)}(\sigma) + L[\sigma, \mathcal{A}, \mathcal{B}] . \quad (30)$$

In the case of absent  $SU(2) \times U(1)$  scalar and gauge fields, a case already discussed in detail in Ref. 6, the purely “inflaton” effective potential  $U_{\text{eff}}(\varphi) = (f_1 e^{-\alpha\varphi} - M_1)^2 [4\chi_2 (f_2 e^{-2\alpha\varphi} + M_2)]^{-1}$  possesses another infinitely large flat region for large positive values of  $\varphi$ :  $U_{\text{eff}}(\varphi) \simeq U_{(+)} \equiv M_1^2 (4\chi_2 M_2)^{-1}$  and the shape of  $U_{\text{eff}}(\varphi)$  is depicted in Fig.1.

It was found in Ref. 6 that this effective “inflaton” potential with two infinitely large flat regions with vastly different scales upon appropriate choice for the values of the parameters accomodates the description simultaneously of the “early” universe evolution (the inflationary epoch) corresponding to the first flat region for large negative  $\varphi$ -values (the left one on Fig.1), as well as the “late” post-inflationary

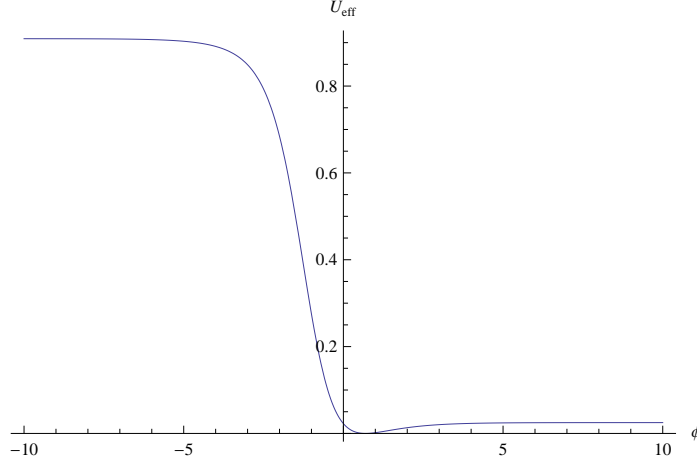


Fig. 1. Qualitative shape of the effective “inflaton” potential  $U_{\text{eff}}(\varphi)$  as function of  $\varphi$ ,  $M_1 > 0$ .

(dark energy dominated) universe corresponding to the second flat region for large positive  $\varphi$ -values (the right one on Fig.1).

The presence of the  $SU(2) \times U(1)$  scalar and gauge fields will not affect significantly the inflationary dynamics in the “early” universe (when  $\varphi$  runs on the left higher flat region of  $U_{\text{eff}}(\varphi)$ , Fig.1), and where the  $SU(2) \times U(1)$  iso-doublet  $\sigma$ -field is massless, see (28)) – this dynamics is governed by the purely kinetic “k-essence” “inflaton” part  $A_{(-)}X + B_{(-)}X^2 - U_{(-)}$  of (27).

However, in the post-inflationary epoch (when  $\varphi$  runs on the right lower flat region of  $U_{\text{eff}}(\varphi)$ , Fig.1) the presence of the  $SU(2) \times U(1)$  iso-doublet  $\sigma$  contributes very significantly to the full effective scalar potential (29), which emerges in an exactly Higgs-like form:

$$U_{(+)}(\sigma) = \frac{1}{4\chi_2 M_2} \left( m_0^2 \sigma_a^* \sigma_a - M_1 \right)^2 + 2\Lambda_0 \equiv \lambda (\sigma_a^* \sigma_a)^2 - \mu^2 \sigma_a^* \sigma_a + \text{const}, \quad (31)$$

$$\lambda = \frac{m_0^4}{4\chi_2 M_2}, \quad \mu^2 = \frac{M_1 m_0^2}{2\chi_2 M_2}. \quad (32)$$

Spontaneous  $SU(2) \times U(1)$  symmetry breakdown occurs at the vacuum value:

$$|\sigma_{\text{vac}}| = \frac{1}{m_0} \sqrt{M_1}, \quad (33)$$

with a small residual cosmological constant  $\Lambda_0$ , which is identified with the current epoch observable cosmological constant. Let us also note that in the post-inflationary epoch according to (30) and (31) the massless “inflaton”  $\varphi$  (entering the kinetic term  $X$  (21) only) loses all interactions with the  $SU(2) \times U(1)$  iso-doublet  $\sigma$  (which has now become the Higgs field), thus avoiding undesirable fifth-force problems.

The dependence of the Higgs-like parameters (32) and the Higgs-like v.e.v. (33) on the integration constants  $M_{1,2}$  and  $\chi_2$  explicitly reveals the nature of the Higgs-like spontaneous gauge symmetry breaking (31) as a gravity-assisted dynamically generated phenomenon in the post-inflationary epoch. It is triggered exclusively through the special non-Riemannian volume-form dynamics in the original gravity-gauge-field-matter action (1) leading to the remarkable two-flat-region shape of the “inflaton” effective potential (Fig.1). Let us stress again that in the “early” universe there is no spontaneous breaking of  $SU(2) \times U(1)$  gauge symmetry and the Higgs-like iso-doublet scalar field  $\sigma$  is massless there (27)-(28).

We conclude with a note on the plausible numerical values for some of the parameters involved. In Ref. 6 it was argued that it is natural to associate  $M_1 \sim M_{EW}^4$ ,  $M_2 \sim M_{Pl}^4$  and  $\chi_2 \sim 10^{-1}$ , where  $M_{EW}$  and  $M_{Pl}$  are the electroweak and Planck scales, respectively. Then, from (33) we see that it is natural also to associate the bare mass of the  $SU(2) \times U(1)$  iso-doublet  $\sigma$ -field  $m_0 \sim M_{EW}$ , so that we will have for the Higgs-like v.e.v.  $|\sigma_{vac}| \sim M_{EW}$  conforming to the standard electroweak phenomenology.

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### Appendix A. Canonical Hamiltonian Treatment of Gravity-Matter Theories with Non-Riemannian Volume-Forms

Let us briefly discuss the application of the canonical Hamiltonian formalism to the gravity-matter model based on two non-Riemannian spacetime volume-forms (1). For convenience we introduce the following short-hand notations for the field-strengths (2), (10) of the auxiliary 3-index antisymmetric gauge fields (the dot indicating time-derivative):

$$\Phi_1(A) = \dot{A} + \partial_i A^i \quad , \quad A = \frac{1}{3!} \varepsilon^{ijk} A_{ijk} \quad , \quad A^i = -\frac{1}{2} \varepsilon^{ijk} A_{0jk} \quad , \quad (\text{A.1})$$

$$\Phi_2(B) = \dot{B} + \partial_i B^i \quad , \quad B = \frac{1}{3!} \varepsilon^{ijk} B_{ijk} \quad , \quad B^i = -\frac{1}{2} \varepsilon^{ijk} B_{0jk} \quad , \quad (\text{A.2})$$

$$\Phi(H) = \dot{H} + \partial_i H^i \quad , \quad H = \frac{1}{3!} \varepsilon^{ijk} H_{ijk} \quad , \quad H^i = -\frac{1}{2} \varepsilon^{ijk} H_{0jk} \quad , \quad (\text{A.3})$$



The canonical momenta conjugated to (A.1)-(A.3) read:

$$\pi_A = R + L^{(1)} - \frac{4\Lambda_0}{\sqrt{-g}} \left( \dot{A} + \partial_i A^i \right), \quad (\text{A.4})$$

$$\pi_B = L^{(2)} + \frac{1}{\sqrt{-g}} \left( \dot{H} + \partial_i H^i \right), \quad \pi_H = \frac{1}{\sqrt{-g}} \left( \dot{B} + \partial_i B^i \right), \quad (\text{A.5})$$

and:

$$\pi_{A^i} = 0, \quad \pi_{B^i} = 0, \quad \pi_{H^i} = 0. \quad (\text{A.6})$$

The latter imply that  $A^i, B^i, H^i$  are in fact Lagrange multipliers for certain first-class Hamiltonian constraints (see Eqs.(A.8)-(A.9) below).

Using the short-hand notation  $(u, \dot{u})$  to collectively denote the set of the basic gravity-matter canonical variables  $(u) = (g_{\mu\nu}, \varphi, \sigma, \mathcal{A}_\mu^A, \mathcal{B}_m)$  and their respective velocities, we have for their respective canonically conjugated momenta:

$$p_u = \left( \dot{A} + \partial_i A^i \right) \frac{\partial}{\partial \dot{u}} \left( R + L^{(1)} \right) + \left( \dot{B} + \partial_i B^i \right) \frac{\partial}{\partial \dot{u}} L^{(2)} \quad (\text{A.7})$$

Now, using relations (A.4)-(A.7) we arrive at the following Dirac-constrained canonical Hamiltonian of the model (1):

$$\begin{aligned} \mathcal{H}_{\text{can}} = & p_u \dot{u} - \frac{\sqrt{-g}}{8\Lambda_0} \left( R + L^{(1)} - \pi_A \right)^2 - \pi_H \sqrt{-g} L^{(2)} \\ & + \sqrt{-g} \pi_H \pi_B - \partial_i A^i \pi_A - \partial_i B^i \pi_B - \partial_i H^i \pi_H. \end{aligned} \quad (\text{A.8})$$

The last three terms in (A.8) show that, indeed, the auxiliary gauge field components  $A^i, B^i, H^i$  (with vanishing canonically conjugated momenta (A.6)) are Lagrange multipliers for the Dirac first class constraints:

$$\pi_A = -M_1 = \text{const}, \quad \pi_B = -M_2 = \text{const}, \quad \pi_H = \chi_2 = \text{const}; \quad (\text{A.9})$$

which are the canonical Hamiltonian counterparts of Lagrangian constraint equations of motion (15).

To conclude, the canonical Hamiltonian treatment of (1) reveals the (almost) pure-gauge nature of the auxiliary 3-index antisymmetric tensor gauge fields  $A_{\mu\nu\lambda}, B_{\mu\nu\lambda}, H_{\mu\nu\lambda}$  – building blocks of the non-Riemannian spacetime volume-form formulation of the modified gravity-matter model (1). Namely, the canonical momenta  $\pi_A, \pi_B, \pi_H$  conjugated to the “magnetic” parts  $A, B, H$  (A.1)-(A.3) of the auxiliary 3-index antisymmetric tensor gauge fields are constrained through Dirac first-class constraints (A.9) to be constants identified with the arbitrary integration constants  $\chi_2, M_1, M_2$  (15) arising within the Lagrangian formulation of the model. The canonical momenta  $\pi_A^i, \pi_B^i, \pi_H^i$  conjugated to the “electric” parts  $A^i, B^i, H^i$  (A.1)-(A.3) of the auxiliary 3-index antisymmetric tensor gauge field are vanishing (A.6), which makes the latter canonical Lagrange multipliers for the above Dirac first-class constraints.

Thus, the method of employing non-Riemannian spacetime volume-forms in gravity-matter theories does not introduce any additional propagating field-theoretic degrees of freedom apart from the standard ones.

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## References

1. J. Bekenstein, *Found. Phys.* **16**, 409 (1986).
2. P. Moniz, P. Crawford and A. Barroso, *Class. Quantum Grav.* **7** (1990) L143.
3. E. Guendelman, *Mod. Phys. Lett.* **A14** (1999) 1043-1052 (*arXiv:9901017*).
4. E. Guendelman and A. Kaganovich, *Phys. Rev.* **D60** (1999) 065004 (*arXiv:gr-qc/9905029*).
5. E. Guendelman, E. Nissimov, S. Pacheva and M. Vasilhoun, in *Eight Mathematical Physics Meeting*, ed. by B. Dragovic and I. Salom (Belgrade Inst. Phys. Press, Belgrade, 2015), pp.105-115. (*arXiv:1501.05518*)
6. E. Guendelman, R. Herrera, P. Labrana, E. Nissimov and S. Pacheva, *Gen. Rel. Grav.* **47**, art.10 (2015). (*arXiv:1408.5344v4*)
7. E. Guendelman, E. Nissimov and S. Pacheva, *Eur. Phys. J.* **C75**, 472-479 (2015) (*arXiv:1508.02008*).
8. E. Guendelman, E. Nissimov and S. Pacheva, *Eur. Phys. J.* **C76**:90 (2016) (*arXiv:1511.07071*).
9. E. Guendelman, E. Nissimov and S. Pacheva, *Int. J. Mod. Phys.* **A30** (2015) 1550133 (*arXiv:1504.01031*).
10. T. Chiba, T. Okabe and M. Yamaguchi, *Phys. Rev.* **D62** (2000) 023511 (*arXiv:astro-ph/9912463*).
11. C. Armendariz-Picon, V. Mukhanov and P. Steinhardt, *Phys. Rev. Lett.* **85** (2000) 4438 (*arXiv:astro-ph/0004134*).
12. C. Armendariz-Picon, V. Mukhanov and P. Steinhardt, *Phys. Rev.* **D63** (2001) 103510 (*arXiv:astro-ph/0006373*).
13. T. Chiba, *Phys. Rev.* **D66** (2002) 063514 (*arXiv:astro-ph/0206298*).